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NASA TECHNICAL TRANSLATION NASA TT F-14,673

INVESTIGATION OF PLASMA HEATING IN A STRONG VARIABLE ELECTRIC FIELD

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Translation of: "Issledovaniye nagreva plasmy v sil'nom peremennom elektricheskom pole," Khar'kov, Physical-Technical Institute, Academy of Sciences Ukrainian SSR, Report KhFTI-70-64, 1970, 41 pages.

NATIONAL AERONAUTICS AND SPACE ADMINISTRATION
WASHINGTON, D. C. 20546 FEBRUARY 1973

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STANDARD TITLE PAGE

1. Report No.	2. Government Accession No.	3. Recipient's Catalog No.
TT-F-14,673		
4. Title and Subtitle INVESTIGATION OF PLASMA HEATING IN A STRONG VARIABLE ELECTRIC FIELD		5. Report Date February, 1973 6. Performing Organization Code
7. Author(s) L. I. Grigoryeva, V. L. Sizonenko, B. I. Smerdov, K. N. Stepanov, and V. V. Chechklin		8. Performing Organization Report No. 10. Work Unit No.
9. Performing Organization Name and Address SCITRAN, P. O. Box 5456 Santa Barbara, California 93108		11. Contract or Grant No. NASW-2483 13. Type of Report and Period Covered Translation
12. Sponsoring Agency Name and Address National Aeronautics and Space Administration, Washington, D. C., 20546		14. Sponsoring Agency Code

15. Supplementary Notes

Translation of "Issledovaniye Nagreva plasmy v sil'nom peremonnom elektricheskom pole," Khar'kov, Physical-Technical Institute, Academy of Sciences Ukrainian SSR, Report KhFTI-70-64, 1970, 41 pages.

16. Abstract

Scattering of charged particles by turbulent pulsation of an electric field leads to effective heating of plasma. The experiments showed that, during turbulent heating of plasma by fast magneto-acoustic waves, preferential heating of ions was correlated with small amplitudes, while at large amplitudes both electrons and ions were heated. The plasma was created by 18-microsecond pulse discharges with oscillating electrons in hydrogen. Measurement of the longitudinal energies of electrons and ions emerging from the plasma along a fixed magnetic field was accomplished with the aid of a three-network analyzer with retarding potential. Observations were also made on the instability of plasma with a transverse current, heating of electrons in the field of a turbulent plasma, and ion heating.

17. Key Words (Selected by Author(s))	18. Distribution Statement
	Unclassified - Unlimited

19. Security Classif. (of this report)	20. Security Classif. (of this page)	21. No. of Pages	22. Price
Unclassified	Unclassified	~~41~~ 43	

NASA-HQ

INVESTIGATION OF PLASMA HEATING IN A STRONG VARIABLE ELECTRIC FIELD

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I. Introduction /3*

As is known, when an electric current passes perpendicularly to an external magnetic field, excitation of different oscillations is possible in the plasma, if the relative velocity of the electrons and ions $\vec{U} = \vec{U}_i - \vec{U}_e$ exceeds a certain critical value $U_{cz} \sim \sqrt{(T_i + T_e)/m_i}$. Scattering of charged particles by turbulent pulsations of the electric field, which are caused by the development of these instabilities, leads to effective plasma heating. Anomalous absorption of high-frequency energy and plasma heating in strong variable electric fields, which are perpendicular to the external magnetic field, were observed in many experiments (see, for example, [1 — 7]). In these experiments, the effective collision frequency, determining the absorption of the high-frequency field energy and plasma heating, greatly exceeded the frequency of paired Coulomb collisions. The observed absorptions of the high-frequency field energy could not be explained by Cherenkov or cyclotron absorption by electrons and ions in the plasma.

The small scale hydrodynamic instability of the plasma with a transverse current was observed in experiments on plasma heating by a strong magneto-sonic wave (whistle) with a large

*Numbers in the margin indicate pagination of the original foreign text.

amplitude [8, 9] in Θ-pinch [10] and in experiments with plasma bunches entering a nonhomogeneous magnetic field [11]. Small scale instabilities were also observed at the front of a collisionless shock wave (see, for example, [12], et al.).

In these studies, there was predominant heating of either ions or electrons, or the electrons and ions were heated to the same order of magnitude.

The reasons for this are still not clear. Which of the numerous current instabilities was responsible for the observed heating is not clear either. Theoretical studies have shown that in a plasma with a transverse current excitation of different types of oscillations is possible [3, 13 — 18]. The role of each of these instabilities in the problem of anomalous resistance and plasma heating may be only clarified after a nonlinear theory is formulated, whose development was begun in [19 — 25]. /4

It was shown experimentally in this study that in the case of turbulent plasma heating with a rapid magneto-sonic wave (whistle) at low amplitudes primarily ions are heated, and at large amplitudes electrons and ions are heated. In the latter case, the electrons are heated in a very small period of time (in a half period of whistle oscillations $\pi/\Omega \approx 7 \cdot 10^{-8}$ sec). A theoretical examination shows that plasma instability may be responsible for the observed rapid heating of electrons. This instability is due to excitation of "magnetized" ion-sonic oscillations in the case of relative electron and ion motion at a rate of $U \gg U_{cz}$ perpendicularly to the magnetic field. The characteristic frequencies and increase of these oscillations are on the order of

$$\mathrm{Re}\,\omega \sim k v_s , \quad \mathrm{Im}\,\omega \sim \sqrt{\omega_{He}\omega_{Hi}} , \tag{1}$$

where $v_j = \sqrt{T_e/m_i}$ is the velocity of ion sound, the wavelength is less than, or on the order of, the Larmor electron radius $\rho_e = v_{T_e}/\omega_{H_e}$ and the Debye electron radius

$$v_{Te}/u \gtrsim k\rho_e \gtrsim 1 , \qquad k r_D \lesssim 1 ,$$

The direction of propagation is almost perpendicular to the magnetic field /5

$$(m_e/m_i)^{1/2} < \cos\theta < u/v_{Te} < 1 ,$$

The phase velocity of the oscillations along the magnetic field is on the order of the thermal electron velocity ($|\omega|/k_{\parallel} \sim v_{T_e}$).

Ions may be heated by the stochastic mechanism of the interaction of non-resonance ions with ion-sonic and hydrodynamic oscillations, caused by the finite nature of the time of correlation of turbulent pulsations of the electric field [26]. For hydrodynamic oscillations, the frequency and increase are on the order of $\sqrt{\omega_{H_e}\omega_{H_i}}$, $k\rho_e \ll 1$ and $\cos\theta \sim (m_e/m_i)^{1/2}$. This mechanism is important at small values of the current velocity and for heating the plasma electron component.

2. Experimental Arrangement. Measurement Method

A detailed description of the experimental arrangement is given in [7]. The plasma was produced by pulsed discharge with oscillating electrons in hydrogen at pressures of $\lesssim 10^{-3}$ mm Hg. The discharge current lasted 18 microseconds. The inner diameter of the glass discharge tube was $2\alpha = 6.6$ cm; the distance between the cathodes was 88 cm. The quasi-constant magnetic field H_0 with a corkscrew geometry (corkscrew ratio — 1.4) had a homogeneous section in the middle section at a length of 70 cm.

Depending on the initial hydrogen pressure and the magnetic field strength, the average (over the radius) density of the plasma $\bar{n}$, produced in the discharge, was measured between 10^{13} — 10^{14} cm^{-3}. The distribution of the plasma density over the radius $n_o(\tau)$, measured by means of a double Langmuir probe, was well approximated by a linear function $n_o(\tau) = 2\bar{n}(1 - \tau/a)$ at every moment of time after the discharge current was stopped. /6

The source of the high-frequency energy was a pulsing circuit with an energy up to 10 J. The inductance of the circuit was comprised of eight different sections encompassing the discharge tube and connected to the circuit condensors in pairs in opposite phase. The axial period of the electromagnetic field produced by a coil was Λ = 20 cm. The circuit oscillation frequency was $\Omega/2\pi$ = 7 MHz, and the quality without the plasma was 35.

As was shown in [7], in excitation resonance of a rapid magneto-sonic wave, i.e., when the longitudinal length of a rapid magneto-sonic wave in a plasma coincides with Λ, more than fifty percent of the energy in the circuit is expended on the plasma, and there is rapid ($\lesssim$ 1 μsec) plasma heating (primarily of electrons) up to a temperature on the order of 100 eV. The variable magnetic field on the system axis increased in resonance by approximately a factor of two. As it was shown in [27], in the region of excitation resonance of the plasma, a complex spectrum of oscillations arose, due to the non-stationary nature of the process and the phenomena of nonlinear interaction of the waves.

In this study, we measured the average (over the radius) density of the electrons by means of a microwave interferometer. As was shown above, the local density may be determined by means of a double Langmuir probe, which is introduced into the plasma from the end of the discharge tube, and may move along the axis of the system and along the radius. A similar method may be used in a discharge to introduce a miniature high frequency magnetic probe, by means of which the variable field $\tilde{H}_z$ on the axis of the system may be measured.

The longitudinal energies of the electrons and ions leaving /7 the plasma along a constant magnetic field were measured by means of a three-grid analyzer with a retarding potential [28]. This analyzer was placed on the system axis behind the magnetic mirror. An opening was made in the cathode so that the plasma could enter the analyzer. During the measurement process, the dependence of the current of ions or electrons at the collector upon the retarding potential was determined. Using this dependence, by differentiation we may find the function of the longitudinal velocity distribution of particles entering the analyzer, and then the average longitudinal energy of the particles. These data may be used to determine the particle heating in the plasma.

The density of the plasma transverse energy was determined from the diamagnetic signal. The diamagnetic probe was a multi-loop coil in an electrostatic screen on the discharge tube between two adjacent sections of the circuit coil, at the nodal point of the field $\tilde{H}_z$ produced by the coil.

The effectiveness with which the high frequency energy was transferred from the circuit into the plasma is characterized by the transition coefficient α, which represents the ratio of the energy, absorbed in the circuit in the presence of a plasma with subtraction of the ohmic losses in the circuit itself, to the total energy supplied to the circuit condensors before it is turned on. The quantity α is determined by the device for measuring the damping decrement [29]. A high frequency magnetic probe is the sensor for this device. This probe produces a signal which is proportional to the current in the circuit.

After the experimental results given below are obtained, the pulsing circuit is turned on for several microseconds after the discharge current has ended, at the moment when the average (over the radius) density of the electrons of the disintegrating plasma decreases to $\bar{n} \approx 3.10^{13}$ cm^{-3}. The increase in the plasma density due to additional ionization after the circuit is switched on does not exceed ~ 20%. The capacitance of the circuit is charged to a voltage of 54 kV, which corresponds to an energy of 7.2 J. Thus, the maximum amplitude of the field $\tilde{H}_z$ on the axis under the coil section is ~ 250 Oe in a vacuum and ~ 500 Oe in a plasma at the excitation resonance ($H_0 \approx 800$ Oe). /8

Figure 1 shows oscillograms of several quantities characterizing plasma heating by a high frequency field. The measurements were performed in the excitation resonance region. It follows from Figure 1 that, after the current is turned on in the circuit (Figure 1a), the axial variable magnetic field $\tilde{H}_z$ in the plasma (Figure 1b) begins to increase. The field amplitude in the second half-period of oscillations is much greater than in the first half-period, i.e., there is resonance buildup of oscillations in the plasma. When the high frequency field reaches a maximum in the second half-period of oscillations, its amplitude sharply decreases, reaching only ~ 1/4 of the maximum value in the third half-period. Simultaneously with strong damping of the high frequency field, there is a rapid increase in the diamagnetic signal (Figure 1c). First, the high frequency field produced by the circuit is superimposed on the diamagnetic probe. This interference assists in establishing the time of increase of the diamagnetic signal, which is approximately

half of the oscillation period, i.e., 0.07 μsec. An increase in the diamagnetism is only caused by plasma heating, since in such a short period of time there cannot be any great increase in the plasma density due to additional ionization caused /9
by electron heating. At subsequent moments of time, along with oscillations at the frequency Ω, there are also oscillations in the plasma at other frequencies, which are visible not only on oscillograms of the field $\tilde{H}_z$, but also on the diamagnetic signal. These oscillations were studied in detail in [27]. It is interesting to note that, after a sharp increase in the plasma temperature and a decrease in the amplitude of oscillations at a frequency of Ω, these oscillations are damped to a much lesser extent, as follows from Figure 1b.

Rapid damping of the oscillations in a plasma, accompanied by its heating, may be due to a sharp increase in the effective frequency of particle collisions in a strong high frequency field. The possible reasons for such an increase in the effective frequency of the collisions will be examined in Sections 4 — 7.

The average (over the radius) maximum plasma temperature, determined from the signal of diamagnetism with allowance for the plasma density distribution over the radius, was $T_e + T_i \approx 150$ eV for the case under consideration. In order to clarify the manner in which the ion and electron temperatures change in particular, let us turn to Figures 1b and 1e, which show oscillograms of ion currents (Figure 1b) and electron currents (Figure 1e) on the analyzer collector with a zero retarding potential. It may be seen from these figures that, after the high frequency field is switched on, there is a rapid increase in the ion and electron currents. This increase in the currents, just like an increase in the diamagnetic signal, is primarily caused by heating of the electrons and ions. It is thus apparent that the electrons group simultaneously with a sharp decrease in the amplitude of the high frequency field in the plasma and an increase in the diamagnetic signal. The ion current reaches a maximum value only ~ 1 μsec

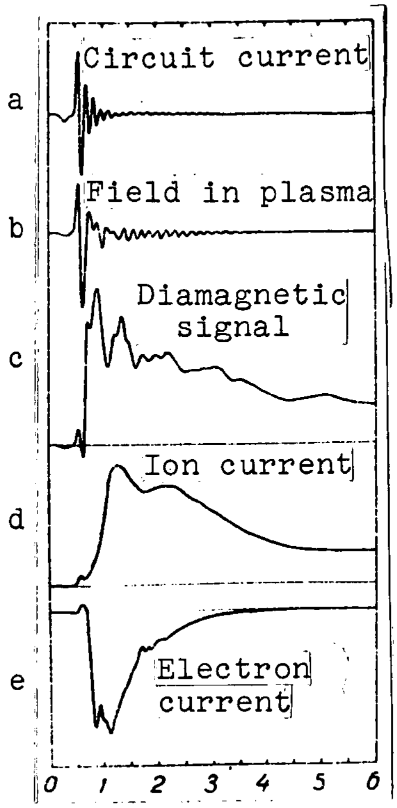

Figure 1.

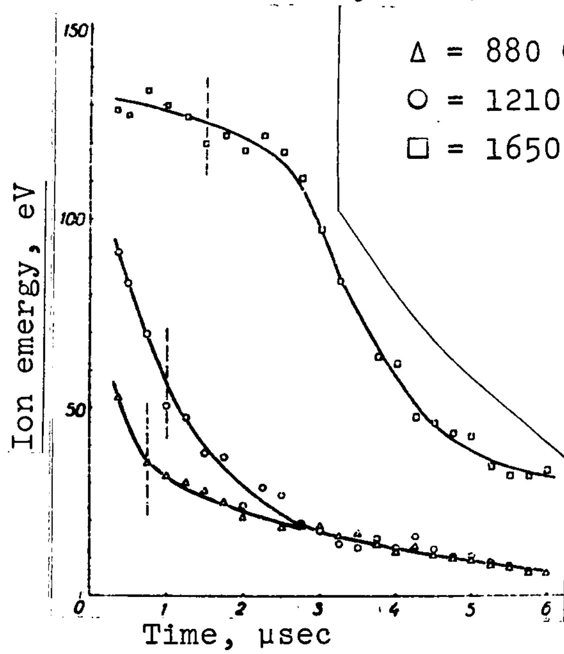

Figure 2.

/10

after the high frequency field is turned on, or ~ 0.5 μsec after the beginning of electron heating. However, it is impossible to determine the ion heating time precisely, /11

since the retardation time of the ion current maximum is influenced not only by the finiteness of the heating time, but also by the finiteness of the time of flight of an ion from the heating location to the collector (the corresponding distance is not known exactly, and may be on the order of several centimeters which, for an ion velocity of $\lesssim 10^7$ cm/sec occurs in the conditions under consideration, and this may produce the observed retardation time). To determine the ion temperature which is reached under these conditions we shall use the fact that at later moments of time, when the high frequency field in the plasma is small, the electron current on the analyzer collector decreases more rapidly than the ion current, which points to more rapid cooling of the electrons. Comparing Figures 1c, 1d, and 1e, we may assume that 2 — 3 microseconds after the high frequency field is switched on, the diamagnetism of the plasma is primarily caused by ions. Thus the ion temperature determined from the diamagnetic signal equals 30 — 40 eV (under the assumption that the plasma density is distributed evenly over the radius and is 3.10^{13} cm^{-3}). This value agrees satisfactorily with the results given below from direct measurements of ion energy by the retarding potential method.

It must be noted that disconnecting the magnetic mirrors does not lead to any important change in the oscillograms shown in Figure 1, even at the initial moments of time when the plasma temperature exceeds 100 eV. Just as in the case of rapid heating, this insensitivity of a hot plasma to magnetic mirrors is apparently due to the anomalously high effective frequency of particle collisions in a strong high frequency field.

Let us now examine in greater detail the experimental data /12
characterizing the heating of ions by a high frequency field. These data are based on measuring the dependence of the ion current at the collector of a multi-grid analyzer on the retarding potential at different moments of time after the high frequency field is turned on and for various strengths of a constant magnetic field. After processing the curves on a computer, the average potential energy of ions entering the analyzer was found. As is known [35], the longitudinal energy distribution of ions entering the input of an analyzer is generally speaking not identically distributed in the plasma. As was indicated above,

the effect of a finite flight time of the ions may have a substantial influence on the distribution within the analyzer. This influence must be particularly great at the initial moments of time. The dependence of the average ion energy on the time after a high frequency field is switched on for three different values of H_o is shown in Figure 2. The moments of time when the current of ions at the collector of the analyzer reaches a maximum value are shown on the figure by the vertical dashed lines. It is natural to assume that the average energy of ions entering the analyzer, which is determined at this moment, is closest to the greatest average energy which the ions have in a plasma during heating. The average ion energy determined in this way actually coincides in order of magnitude with the ion energy in the plasma determined from the diamagnetic signal. The most noteworthy factor pertaining to ion heating is the fact that the maximum ion energies as a function of H_o are not reached for the resonance value of H_o, where the largest values of the high frequency energy /13
absorbed by a plasma, the amplitude of a high frequency field in the plasma (Figure 3a), and the electron energy (Figure 3b) are assumed. As follows from Figure 3b, the ion energy increases with an increase in H_o, and assumes the largest value (130 eV at the ion current maximum) for H_o which exceeds the resonance value by a factor of approximately two. At the values of the constant magnetic field strength under consideration, the Larmor radius of ions heated by a high frequency field may be compared with the radius of the discharge tube. The upper limit of the ion energy is apparently determined by this factor. In actuality, as follows from Figure 3c, for all values of H_o considered, the Larmor radius of an ion, corresponding to the average energy shown in Figure 3b, changes very little, remaining at the level $\sim (0.3 - 0.2)a \sim 1$ cm.

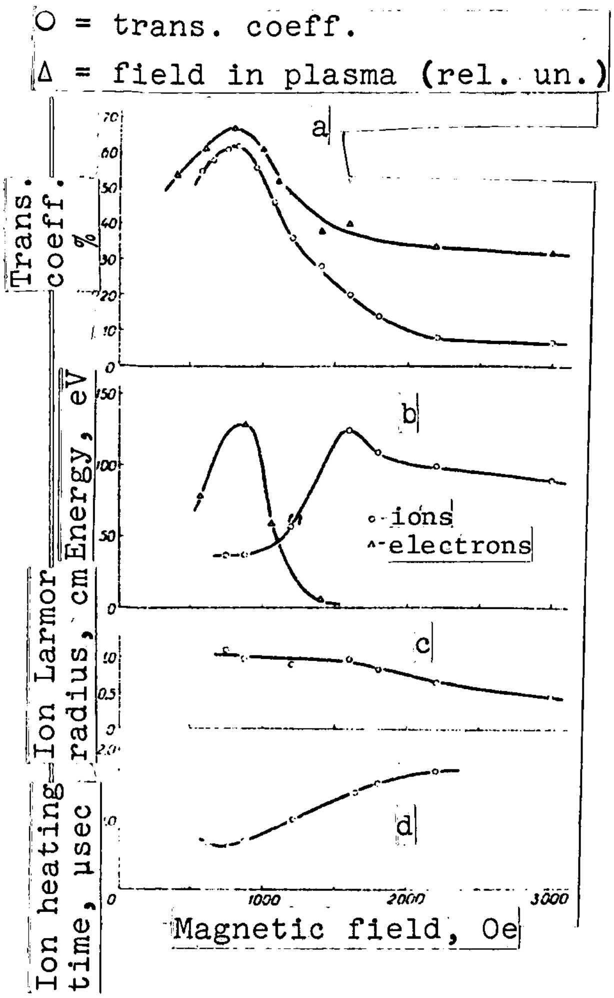

Figure 3.

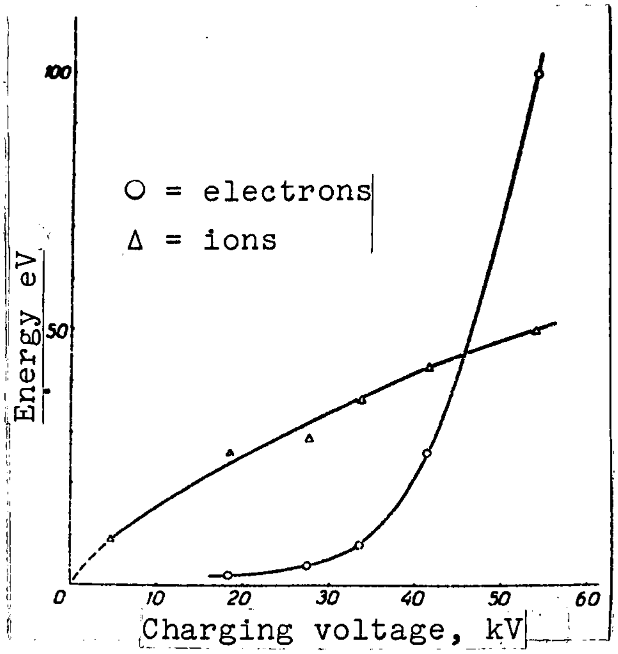

Figure 4.

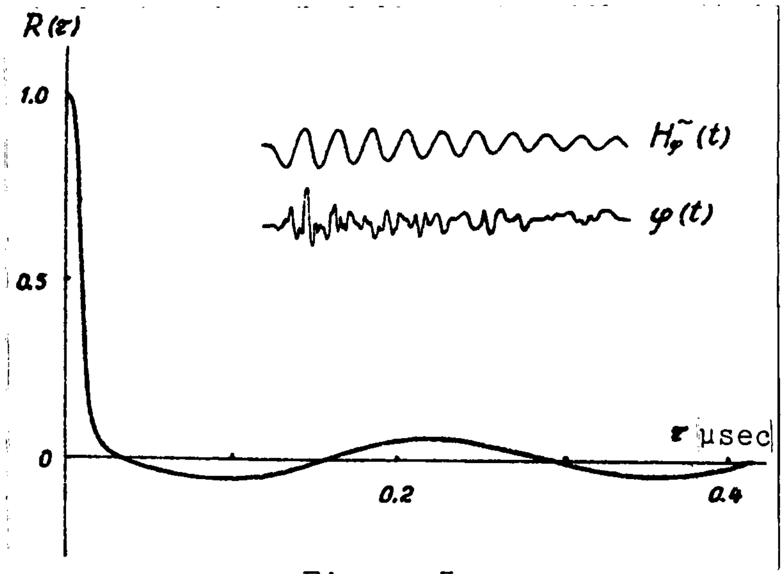

Figure 5.

Figure 3 also shows the dependence on H_o of the time at which the ion current maximum appears with respect to the moment at which the high frequency field is turned on with a zero retarding potential (Figure 3d). It follows from Figure 3d that this time increases as the distance increases from excitation resonance toward large H_o (apparently this will occur also for a decrease in H_o). Apparently, this is caused by a decrease in the ion heating rate, which in its turn is caused by a decrease in the amplitude of high frequency oscillations in the plasma (see Figure 3a). As follows from Figure 3c, the ions acquire a great energy in the high frequency field. Hence, with departure from resonance, the damping coefficient of oscillations in the circuit decreases, and therefore the effective lifetime of a high frequency pulse increases.

As is shown by measurements (Figure 3b), the energy of electrons, which is the greatest at excitation resonance, rapidly decreases with departure from resonance to a value on the order of several eV. Thus, for a large H_o, practically all of the high /15 frequency energy absorbed by the plasma ($\alpha \sim 10\%$) goes into heating the ions. It may be assumed that the rate at which the electrons are heated, and consequently the maximum electron temperature, depend greatly on the amplitude of the high frequency field in the plasma. This assumption (just like the assumption given above regarding the dependence of the ion heating rate on the amplitude of the high frequency field) is confirmed by results of direct measurements of the dependence of ion and electron energy on the charged voltage in the circuit with a fixed value of H_o = 900 Oe, which somewhat exceeds the resonance value (Figure 4). Under these conditions, the amplitude of the high frequency field in the plasma in the first period of oscillations is proportional to the charged voltage. As would be expected, with a decrease in the amplitude of the high frequency field in the plasma the electron energy decreases much more rapidly than that of the ions.

In the light of these data, it is understandable that under excitation resonance conditions (Figure 1) significant electron heating occurs only after the second half-period of oscillations: the amplitude of the high frequency field in the first half-period is not large enough to provide a high electron heating rate. As follows from Figure 1b, in the second half-period the amplitude of the high frequency field exceeds the amplitude in the first half-period by a factor of 1.5. It follows from Figure 4 that the corresponding electron energies may differ by one order of magnitude.

4. Instability of a Plasma with a Transverse Current

Let us examine the excitation of small-scale high frequency longitudinal oscillations in a plasma with a transverse current, which may be responsible for turbulent plasma heating in our experiments. We shall assume that the frequency of electromagnetic oscillations (in our case "whistles") caused by the motion of electrons and ions perpendicularly to the magnetic field is much less than the frequency and increment of increase of longitudinal oscillations, and the wavelength is much greater than the wave- /16
lengths of longitudinal oscillations. In this case, during the development of oscillations we may disregard the dependence of electron velocity (and ion) $\vec{U}_{e,\ i}(\vec{r},\ t)$ on the coordinates and time (adiabatic approximation). The increment of increase and the frequency of longitudinal oscillations are much greater than the ion cyclotron frequency, and the wavelength is much less than the Larmor radius of ions, so that the action of the magnetic field on the movement of ions during the development of instability may be disregarded. Then the complex frequency of oscillations with a phase velocity, which is much greater than the thermal velocity of ions and the speed of sound, is determined by the following expression [15]:

$$\omega = \vec{K}\vec{u}_i \pm \delta\omega , \qquad \delta\omega = \frac{\omega_{pi}}{\sqrt{1+\delta\varepsilon_e(\vec{K}, \vec{K}\vec{u}_i)}} , \tag{2}$$

where $\delta\varepsilon_e(\vec{K}, \omega)$ is the contribution to the longitudinal dielectric constant of the plasma electrons. In the case of a Maxwell velocity distribution of electrons

$$\delta\varepsilon_e = \frac{\omega_{pe}^2}{K^2 V_{Te}^2}\left[1 + i\sqrt{\pi}\, z_0 \sum_{n=-\infty}^{\infty} A_n(x)\, w(z_n)\right] , \tag{3}$$

where

$$z_n = \frac{\omega - n\omega_{He} - \vec{K}\vec{u}_e}{\sqrt{2}\, K_{\parallel} V_{Te}} , \qquad w(z) = e^{-z^2}\left(1 + \frac{2i}{\sqrt{\pi}}\int_0^z e^{t^2} dt\right) ,$$

$$A_n(x) = e^{-x} I_n(x) , \quad x = K_\perp^2 \rho_e^2 = \left(\frac{K_\perp V_{Te}}{\omega_{He}}\right)^2 ,$$

$$K_\perp = K \sin\theta , \quad K_\parallel = K \cos\theta , \quad V_{Te}^2 = \frac{T_e}{m_e} .$$

Expression (2) may be used if the angle θ between the wave vector $\vec{K}$ and the magnetic field is not close to $\pi/2$, so that /17

$$\cos^2\theta \gg \frac{m_e}{m_i} , \qquad \frac{|\delta\omega|}{K_\parallel V_{Te}} \ll 1 .$$

Since the quantity $\delta\varepsilon_e$ has an imaginary part which differs from zero due to the presence of a Cherenkov and cyclotron absorption, ω always has a positive imaginary part which differs from zero.

If $\omega_{He} \gg K_{\parallel} v_{Te}$, then the cyclotron absorption of oscillations may be disregarded, and we may disregard in (3) all components with $n \neq 0$ (with the exception of the case when $\vec{K}\vec{U} \approx n\omega_{He}$). In this case, for oscillations with $\cos\Theta \sim U/v_{Te}$ the phase velocity of oscillations along the magnetic field is on the order of the thermal velocity of electrons

$$z_0 = \frac{\vec{K}\vec{u}}{\sqrt{2}K_{\parallel}v_{Te}} \sim 1 \quad ,$$

i.e., in the case under consideration, almost all of the electrons (with $|v_{\parallel}| \lesssim v_{Te}$) are resonance electrons and cause a buildup of oscillations.

For a dense plasma ($\omega_{pe} \gg \omega_{He}$) in the case $\cos\Theta \sim U/v_{Te} < 1$ oscillations with a wavelength which is less than the order of magnitude of the Larmor electron radius, but greater than the Debye radius, have an increment of increase which equals (1) in order of magnitude. We should note that this may be used only when the condition $\omega_{He}/K_{\parallel}v_{Te} \gg 1$ is satisfied, i.e.,

$$K\rho_e \ll \frac{v_{Te}}{u} \quad , \tag{4}$$

when the effects of cyclotron absorption are unimportant. In addition, to satisfy the condition $|\delta\omega/K| \gg v_{Ti}$, it is necessary that the plasma be greatly non-isothermal, i.e., $T_e \gg T_i$. /18

Let us determine the level of turbulent noise which is established as a result of the development of oscillations with an increment of increase in a nonlinear mode. We shall assume that stabilization of the oscillations is due to nonlinear interaction of the waves with participation of the electrons. Nonlinear electron oscillations may be described by a kinetic equation for the oscillating part of the electron distribution function $f^{\sim} = f - f_o^{(e)}$:

$$\frac{\partial f^{\sim}}{\partial t} + \vec{v}\frac{\partial f^{\sim}}{\partial \vec{z}} - \frac{e}{m_e c}[\vec{v}\vec{H}_o]\frac{\partial f^{\sim}}{\partial \vec{v}} + \frac{e}{m_e}\nabla\varphi\frac{\partial f_o^{(e)}}{\partial \vec{v}} + \frac{e}{m_e}\left(\nabla\varphi\frac{\partial f^{\sim}}{\partial \vec{v}} - \left\langle\nabla\varphi\frac{\partial f^{\sim}}{\partial \vec{v}}\right\rangle\right) = 0, \quad (5)$$

where the symbol <...> designates static averaging. The nonlinear term

$$\left|\frac{e}{m_c}\nabla\varphi\frac{\partial f^{\sim}}{\partial \vec{v}}\right| \sim \left|\frac{e K_\perp \varphi f^{\sim}}{m_e v_{Te}}\right|$$

in a very nonlinear mode becomes on the order of the linear term $\partial f^{\sim}/\partial t$, which is equal in order of magnitude to $\vec{K}\vec{U}f^{\sim}$ (in the frame of reference where the average electron velocity equals zero) and which is responsible for Cherenkov absorption of oscillations when

$$E^{\sim} = -\nabla\varphi \sim K_\perp\varphi \sim (m_e/e)\,u\,v_{Te}K_\perp \sim (u/c)H_o K\rho_e, \quad (6)$$

where $K \approx K_\perp$ is the characteristic value of the wave vector. The maximum value of the wave vector of unstable oscillations being considered is limited by the condition (4). When the inequality which is opposite to (4) is satisfied, the oscillation increment of increase is much less than $\sqrt{\omega_{He}\omega_{Hi}}$ (only the narrow region of resonance values K for which $KU \approx n\omega_{He}$ may be an exception, see below). Therefore in (6) we limit the values of K to the maximum value $K_{max} \sim (1/\rho_e)(v_{Te}/U)$, and we shall assume that $K_{max}\tau_D \lesssim 1$, i.e., $\omega_{He}v_{Te}/\omega_{pe}U \lesssim 1$. We should note that when $K \sim K_{max}$ the cyclotron excitation of oscillations becomes great, in addition to Cherenkov excitation /19

$$W \sim \frac{1}{8\pi}\left|\omega \frac{d(\delta\varepsilon_e + \delta\varepsilon_i)}{d\omega}\right| \tilde{E}^2 \sim n_0 m_e u^2 K^2 \rho_e^2 . \tag{7}$$

For pulsations of electron and ion density $\tilde{n}_{e,i}$ and an average electron and ion velocity of $\tilde{U}_{e,i}$, we obtain

$$\begin{aligned} \tilde{n}_e \sim \tilde{n}_i &\sim n_0 (u/v_{Te}), \\ \tilde{u}_e \sim u^2/v_{Te}, \quad & \tilde{u}_i \sim (m_e/m_i)^{1/2} u . \end{aligned} \tag{8}$$

The estimates obtained above for the level of turbulent pulsations may also be obtained by using the kinetic equation for waves

$$\frac{\partial W_{\vec{k}}}{\partial t} = 2\gamma_{\vec{k}} W_{\vec{k}} - \int d\vec{k}\, U_{\vec{k},\vec{k}'} W_{\vec{k}} W_{\vec{k}'} , \tag{9}$$

where $W_{\vec{K}}$ is the spectral density of oscillation energy. The expressions for the matrix element $U_{\vec{K}_1\vec{K}'}$ are given in [30, 31] (in the expressions for $U_{\vec{K}_1\vec{K}^1}$ in [30, 31] it is only necessary to replace the frequency $W_{\vec{K}}$ by $\vec{K}\vec{U}$). Taking only the Cherenkov terms into account in $U_{\vec{K}_1\vec{K}^1}$, we obtain in the usual manner the following

$$U_{\vec{K},\vec{K}'} \sim \frac{\sqrt{\omega_{He}\omega_{Hi}}}{n_0 m_e u^2} \frac{1}{K^2 \rho_e^2} \qquad \left(\cos\theta \sim \cos\theta' \sim \frac{u}{v_{Te}}\right).$$

It thus follows that the nonlinear term in (9) becomes on the order of the term responsible for the oscillation buildup exactly at the level of oscillations (6) — (8). /20

Since turbulence is strong in this case, utilizing Equation (9) may only give the correct order of magnitude. However, an estimate based on the use of the theory of weak turbulence makes sense, since it shows that at a level of oscillations less than (6) instability stabilization due to nonlinear buildup of oscillations by electrons is impossible, since in this case the nonlinear component in (9) represents only a small correction to the linear term $2\gamma_{\vec{K}}W_{\vec{K}}$. (We should also note that, in spite of the strong turbulence, the concept of resonance particles can be retained, since $\vec{K}\vec{U} \sim K_{11}v_{11} \gg \tau m\omega \sim \sqrt{\omega_{He}\omega_{Hi}}$, so that the matrix elements $U_{\vec{K}_1\vec{K}^1}$ will represent integrals containing δ — functions $\delta(\omega_{\vec{K}} - K_{11}v_{11}) = \delta(\vec{K}\vec{U} + K_{11}v_{11})$.)

5. Electron Heating

Let us consider heating of electrons and ions in the field of turbulent pulsations (6). Scattering of resonance electrons by plasma oscillations leads to "a spreading apart" of the background distribution function along the magnetic field. To determine this phenomenon, we shall use the quasilinear equation of particle diffusion by the waves

$$\frac{\partial f_0^{(e)}}{\partial t} = \pi \frac{e^2}{m_e^2} \frac{\partial}{\partial v_{\parallel}} \int d\vec{K}\, K_{\parallel}^2 |\varphi_{\vec{K}}|^2 \frac{\partial f_0^{(e)}}{\partial v_{\parallel}} \delta(\vec{K}\vec{u} + K_{\parallel} v_{\parallel}). \quad (10)$$

(We should note that in this case the quasilinear equation (10) which describes the scattering of electrons by oscillations with the participation of one plasmon, is only valid in order of magnitude, since at the level (6) components indicating the scattering of electrons with the participation of two, three, etc., plasmons must be added to the right part of Equation (10). However, since these components are of the same order of magnitude as those considered in (10), the results obtained from (10) are valid in terms of the order of magnitude). /21

Since the spectral intensity $|\varphi_{\vec{K}}|^2$ differs from zero in a wide range of phase velocity values $\Delta|\omega_{\vec{K}}/K_{\parallel}| \sim v_{Te}$, the opposite action of oscillations on resonance particles leads to diffusion of the distribution function in a wide range of longitudinal velocities $|v_{\parallel}| \sim v_{Te}$. The spectrum of unstable oscillations, according to (2), will expand as the electron heating increases, which will lead to further electron heating.

The characteristic time for a change in the function $f_o^{(e)}$ (heating time) is determined according to (10) and (6) by the following relationship:

$$\frac{1}{\tau_e} \sim \omega_{He} \left(\frac{u}{v_{Te}}\right)^3. \quad (11)$$

This electron heating mechanism leads to an increase in the longitudinal electron temperature in the case $K\rho_e \ll v_{Te}/U$. At least two mechanisms are possible which lead in this situation to an increase of the transverse electron temperature. First, in the case $T_e \gg T_i$ and $U \gg v_j$ excitation of "nonmagnetized" ion-sound oscillations, whose frequency and increment of increase are also determined by Expression (2), is possible. Assuming that for unstable, short-wave ion-sound oscillations $K\tau_e \sim 1$ and $K\rho_e \gg 1$ ($\omega_{pe} \gg \omega_{He}$) and $\cos\theta \sim 1$, we obtain from (2) /22

$$\operatorname{Re}\omega = \vec{K}\vec{u}_i - \frac{K v_s}{\sqrt{1+K^2 r_D^2}}, \qquad \operatorname{Im}\omega = \sqrt{\frac{\pi m_e}{8 m_i}}\,\frac{\vec{K}\vec{u}_e - \operatorname{Re}\omega}{(1+K^2 r_D^2)^{3/2}}.$$

Scattering of the electrons by turbulent ion-sound oscillations leads to isotropization of the distribution function [19, 32], and consequently to an equalization of the longitudinal and transverse temperatures. The isotropization time $\tau_i \sim 1, \nu_j$ where ν_j is the effective frequency of electron scattering determined by the following relationship [19]

$$\nu_j \sim \omega_{pe}\frac{W_j}{n_0 T_e}, \tag{12}$$

where W_j is the energy density of ion-sound oscillations. If the ion-sound instability is stabilized due to nonlinear scattering by resonance electrons, just as in an isotropic plasma, then [24] $W_j \sim (m_e/m_i)n_0T_e$, and temperature anisotropy does not occur ($\tau_i < \tau_e$), if

$$\frac{\omega_{He}}{\omega_{pe}} < \left(\frac{v_{Te}}{u}\right)^3 \frac{W_j}{n_0 T_e} \sim \left(\frac{v_{Te}}{u}\right)^3 \frac{m_e}{m_i}.$$

Scattering of electrons by ion-sound oscillations will also lead to their heating. In this case, the heating time is on the order of $\tau_e^{(j)}$, where

$$\frac{1}{\tau_e^{(j)}} \sim \left(\frac{u}{v_{Te}}\right)^2 \nu_s \sim \omega_{pe} \left(\frac{u}{v_{Te}}\right)^2 \frac{W_s}{n_0 T_e} . \tag{13}$$

In this case, the longitudinal and transverse electron temperatures are increased.

Secondly, for shortwave ion-sound oscillations which are distributed almost perpendicularly to the magnetic field ($\cos\Theta \sim U/v_{Te}$), when condition (4) is disturbed, i.e., in the case $K\rho_e \sim v_{Te}/U$, as was already noted above, scattering of electrons by turbulent oscillations becomes significant under conditions of cyclotron resonance $\omega = K_{\parallel}v_{\parallel} + n\omega_{He}$, which leads to an increase in the transverse electron temperature. According to quasilinear theory, the rate at which the transverse temperature increases is determined by the following relationship /23

$$\frac{dT_\perp/dt}{dT_\parallel/dt} \sim \frac{K_\parallel v_\parallel (\partial f_0^{(e)}/\partial v_\parallel)}{\omega_{He}(\partial f_0^{(e)}/\partial v_\perp)} \sim \frac{K_\parallel v_{Te}}{\omega_{He}} .$$

In the case $K\rho_e \sim v_{Te}/U$ and $\cos\Theta \sim U/v_{Te}$, it follows that the rate at which $T_\perp$ increases is the same as the rate at which $T_\parallel$ increases.

Cyclotron instabilities for which $\vec{K}\vec{U} \approx n\omega_{He}$ may also lead to an increase in the transverse energy. The maximum increment of increase of these oscillations which are excited by resonance electrons is reached in the case $\cos\Theta \sim (U\, v_{Te})^{\tau}$ and $K\rho_e \sim v_{Te}/U$ (but thus $K\tau_D \sim \omega_{He} v_{Te}/\omega_{pe} U \lesssim 1$):

$$\mathrm{Im}\,\omega \sim (v_{Te}/U)\sqrt{\omega_{He}\omega_{Hi}}\ .$$

These oscillations develop in a narrow range of values for the wave vector $\Delta K/K \sim U/v_{Te}$, and are apparently stabilized at a comparatively low level.

Finally, for values of the angle Θ which are close to $\pi/2$, coherent (not connected with dissipative phenomena) excitation of shortwave ($K\rho_e \sim n v_{Te}/U \gg 1$) cyclotron oscillations is possible (see also [14]). Assuming in this case that $|Z_n| \gg 1$, /24
we find that

$$\delta\varepsilon_c = -\,\omega_{pe}^2 \sum_{n=1}^{\infty} \frac{n^2 A_n(x)}{(\omega - \vec{K}\vec{U}_e)^2 - n^2\omega_{He}^2}$$

Thus in the case $\delta\varepsilon_e \sim 1$ we find the following for $\vec{K}\vec{U} > \omega_{He}$

$$\left(\frac{\omega_{pe}^2 v_{Te}^2}{\omega_{He}^2 U^2} \gtrsim K^3\rho_e^3 \gg 1\,,\quad \cos\Theta \ll \frac{1}{K\rho_e}\,,\ \frac{U}{v_{Te}} \right).$$

It must be noted, however, that these oscillations cannot be excited, since their wavelength $\lambda_{\parallel} = 2\pi/K_{\parallel}$ is great due to the smallness of $\cos\Theta$, and the condition $\lambda_{\parallel} \ll \Lambda$ may be disturbed, where Λ is the longitudinal wavelength of oscillations with the frequency Ω producing a transverse current (in our experiments, $\lambda_{\parallel} \gg 3$ cm must hold for these oscillations, when $\Lambda = 20$ cm, so that the condition $\lambda_{\parallel} \ll \Lambda$ may be satisfied).

The presence of turbulent pulsations leads to the occurrence of a turbulent friction force acting on the electrons (and ions). We may find this force by multiplying the quasilinear equation for the background distribution function of electrons by $\vec{v}$ and integrating in the space of velocities. Then, taking the fact into account that $\cos\theta \sim U/v_{Te} \sim 1/K\rho_e$, we obtain the following estimate for the effective collision frequency:*

$$\nu_{eff} \sim \frac{u}{v_{Te}}\,\omega_{He}. \tag{14}$$

6. Ion Heating

Since the phase velocity of the oscillations being considered is much greater than the thermal velocity of ions, ion heating by the field of turbulent pulsations under conditions of Cherenkov resonance is impossible. However, due to strong nonlinear interaction of the waves, which leads to a "reduction" of the phase of the Fourier field components and to attenuation of their time correlation, the mechanism of stochastic ion heating becomes very effective. /25

For these oscillations, if the amplitude of the oscillations of directional ion velocity $U_i^{\sim} \sim (m_e/m_i)^{1/2}U$ is much less than the thermal velocity of ions, then the ion distribution function may be represented in the form of a sum $f_i = f_o^{(i)} + f^{\sim}$, where

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*The use of this expression gives the following estimate

$$\Delta \sim (c/\omega_{pe})(H_o^2/4\pi n_o T_e)^{1/2}.$$

for the width of the front of a collisionless shock wave propagated perpendicularly to the magnetic field.

the oscillating part of the distribution function $f^{\sim}$ is much less than the averaged background distribution function $f_o^{(i)}$. Since $U_i^{\sim} \ll v_{Ti}$, we may disregard the nonlinear scattering of waves by ions. In this case, the background distribution function of ions may be described by the equation [26]:

$$\frac{\partial f_o^{(i)}}{\partial t} = \frac{\partial}{\partial v_i} \mathcal{D}_{ij} \frac{\partial f_o^{(i)}}{\partial v_j}, \tag{15}$$

where the diffusion coefficient is equal in order of magnitude to

$$D_{ij} = \frac{e^2}{m_i^2} \int d\vec{K}\, K_i K_j\, |\varphi_{\vec{K}}|^2 \frac{1/\tau}{(\omega - \vec{K}\vec{V})^2 + 1/\tau^2}, \tag{16}$$

The quantity $\tau(\vec{K}_1\vec{v})$ is the characteristic correlation time of $\varphi_{\vec{K}}(t)$. In the case of strong turbulence, the quantity $1/\tau(\vec{K})$ will be on the order of the nonlinear decrement of damping, i.e., $\tau(\vec{K}) \sim 1/\sqrt{\omega_{He}\omega_{Hi}}$. Taking this fact into account, from (15) and (16) we find that the ion temperature changes in the characteristic time τ_i, determined in order of magnitude by the relationship

$$\frac{1}{\tau_i} \sim \left(\frac{eE^{\sim}}{m_i v_{Ti}}\right)^2 \frac{1}{\tau} \sim \omega_{Hi} \sqrt{\frac{m_e}{m_i}} \left(\frac{u}{v_{Ti}}\right)^2. \tag{17}$$

It thus follows that the ion temperature increases as /26

$$T_i \sim \omega_{Hi} \sqrt{\frac{m_e}{m_i}} m_i \int_0^t u^2(t')dt' \tag{18}$$

In addition to the ion-sound and cyclotron oscillations, in the case $U \gg v_j$, v_{Ti} excitation of hydrodynamic oscillations is also possible, for which thermal motion of electrons and ions is not significant, and for which

$$|\omega - \vec{k}\vec{u}_e| \gg k_{\parallel} v_{Te}, \quad |\omega - \vec{k}\vec{u}_i| \gg k v_{Ti}, \quad k\rho_e \ll 1 .$$

These oscillations are propagated almost perpendicularly to the magnetic field ($\cos^2\theta \sim m_e/m_i$), and have a frequency and increment of increase on the order of

$$\mathrm{Re}\,\omega \sim \mathrm{Im}\,\omega \sim \sqrt{\omega_{He}\omega_{Hi}} \qquad (\omega_{pe} \gtrsim \omega_{He}).$$

Stochastic ion and electron heating at hydrodynamic turbulent pulsations also occurs according to (7).*

The time of stochastic ion heating in the case of scattering by nonmagnetized ion-sound pulsations $\tau_i^{(j)}$ is determined by the relationship

$$\frac{1}{\tau_i^{(s)}} \sim \frac{1}{\tau_c^{(j)}} \frac{W_s}{n_0 T_i},$$

*This was shown by the authors together with V. V. Demchenko.

where $\tau_c^{(j)}$ is the correlation time of ion-sound oscillations. It is assumed that the finiteness of the correlation time is due to the nonlinear interaction of oscillations with participation of the electrons, and we set $1/\tau_c^{(j)} \sim \tau m\omega \sim (U/v_{Te})\omega_{pi}$. Then, with allowance for $W_j/n_0T_e \sim m_e/m_i$, we obtain /27

$$\frac{1}{\tau_i^{(j)}} \sim \omega_{pi} \frac{u}{v_{Te}} \frac{T_e}{T_i} \frac{m_e}{m_i} .$$

Let us digress from the possible influence of electron cyclotron oscillations and nonmagnetized ion-sound on plasma heating, and we shall only consider ion-sound oscillations in the case $\cos\Theta \sim U/v_{Te}$ and hydrodynamic oscillations. For the characteristic of the rate at which electrons and ions are heated, let us introduce the parameter $\varkappa$ which is determined by the equation

$$\varkappa = \frac{\tau_e}{\tau_i} \sim \frac{v_s}{u} \frac{T_e}{T_i} . \tag{19}$$

In the case $\varkappa \ll 1$ electrons are heated more rapidly if $\varkappa \gtrsim 1$, and then ions and electrons are heated at approximately the same rate, so that with prolonged heating the temperatures of the ions and electrons must be of the same order of magnitude when there are no losses.

Let us now discuss the manner in which plasma heating occurs when a high frequency field is turned on with a large amplitude. If the amplitude of the high frequency field is so great that $\varkappa \ll 1$, then at first there is primarily electron heating. As the electron temperature increases (and possibly the amplitude of the current velocity U decreases), the parameter $\varkappa$ increases

and when æ becomes greater than unity, the ion temperature begins to "approach" the electron temperature. With a certain value of the ratio $T_e/T_i > 1$, "nonmagnetized" ion-sound oscillations are disrupted, and a further plasma heating only occurs due to hydrodynamic instability until the thermal ion velocity is not comparable with the current velocity U. In this case, the hydrodynamic instability changes into kinetic instability [16] for /28
which

$$Re\,\omega \sim Im\,\omega \sim \sqrt{\omega_{He}\,\omega_{Hi}}\ , \quad |\omega - \vec{k}\vec{u}| \sim k v_{Ti}\ ,$$

Stochastic heating, connected with "reduction" of the phase, is replaced by Cherenkov quasilinear heating of ions and electrons. This last stage of turbulent heating is continued until the thermal ion velocity is no longer close to the critical value of the current velocity $U = U_{cz}$, below which the plasma is stable in the wave field.

This picture of heating is only valid if we may disregard plasma energy losses. If the energy acquired by electrons from turbulent pulsations is compensated by energy losses in the case of inelastic collisions of electrons with neutral particles, and the energy lost by the ions during inelastic collisions is small, then primarily ion heating occurs (even if æ < 1).

Plasma heating may also be greatly slowed down, if there is anomalously high thermal conductivity of the plasma across the magnetic field, and if this conductivity is caused by scattering of particles by turbulent pulsations in a nonhomogeneous plasma.

7. Conclusion

In conclusion, let us formulate the experimental results obtained, and let us compare them with the theoretical determinations.

In the case of excitation of a rapid magneto-sonic wave in a plasma (electron desnity $\bar{n} \approx 3.10^{13}$ cm^{-3}, the temperature at the moment of wave excitation $T_{e,\,i} \lesssim 1$ eV, the wave frequency $\Omega = 4.4\cdot10^{7}$ sec^{-1}, and the longitudinal and radial wave vectors are equal to $K_{\parallel} = 2\pi/\Lambda \approx 0.3$ cm^{-1}, $K_z \approx 1$ cm^{-1}), the wave amplitude under conditions of resonance excitation is increased in three quarters of a period of the oscillations up to the maximum value of $\tilde{H_z} \approx 500$ Oe. As was shown in [27], the dependence of the magnetic field of a whistle on the coordinates in the plasma is satisfactorily described by the relationships of linear theory (see, for example, [34]), in spite of the fact that the amplitude of the variable magnetic field on the $\tilde{H_z}$ axis is comparable with the magnitude of the external constant magnetic field H_0. The magnitude of the current velocity is approximately equal to the drift velocity of electrons in the azimuthal direction /29

$$u \approx u_e \approx c\,\frac{E_z}{H_0} \sim \frac{\Omega \tilde{H_z}}{K_{\parallel} H_0} J_1(K_z z)\,,$$

where E_z is the radial component of the whistle electric field strength. At the point $z = a/2 = 1.6$ cm, where the plasma density equals the average value, $U \approx 5\cdot10^{7}$ cm/sec. The directional velocity of ions is much less

$$u_i \approx \frac{eE_z}{m_i\Omega} \sim \frac{e\tilde{H_z}}{m_i k_{\parallel} c} J_1(K_z z) \sim 6\cdot10^{6}\ \frac{\text{cm}}{\text{sec}}$$

After the time $\sim 10^{-7}$ sec (approximately three quarters of a period of the high frequency field), after the whistle magnetic field has reached a maximum value, rapid heating begins of electrons whose transverse temperature, measured by a diamagnetic probe, is ~ 100 eV, in the time $\lesssim 7 \cdot 10^{-8}$ sec (approximately half of the high frequency field period). In the same time, the whistle amplitude decreases by a factor of four, so that practically all of the whistle energy stored in the plasma $W_W \sim \tilde{H}_z^2/8\pi \sim 5 \cdot 10^3$ erg/cm^3 changes into the electron energy $n_0 T_e \sim 5 \cdot 10^3$ erg/cm^3.

The whistle phase velocity along the magnetic field /30
$\Omega/K_{\parallel} \sim 1.5 \cdot 10^8$ cm/sec is less than the thermal electron velocity $v_{Te} \approx 4 \cdot 10^8$ cm/sec ($T_e \sim 100$ eV). Therefore, the whistle may be absorbed by plasma electrons under conditions of Cherenkov resonance. However, Cherenkov absorption of the whistle cannot be responsible for the rapid damping of the whistle. The whistle damping time due to Cherenkov absorption equals [33]

$$\frac{1}{\tau_C} = \sqrt{\frac{\pi}{2}} \sin^2\vartheta \, K \rho_e \Omega , \qquad (20)$$

where ϑ is the angle between the whistle wave vector and the magnetic field. In our case $\sin^2 \vartheta \approx 0.9$, $K \approx 1$ cm^{-1}, so that $\tau_C \approx 6 \cdot 10^{-7}$ sec. This value of the damping time is one order of magnitude greater than the observed value.

The rapid electron heating may be explained by electron scattering by turbulent shortwave longitudinal oscillations of the "nonmagnetized" ion sound type, which is propagated almost perpendicularly to the magnetic field ($\cos\Theta \sim U/v_{Te} \sim 1/K\rho_e$). Under the conditions of our experiment, these oscillations have the characteristic frequency $\omega \sim Kv_j \sim (v_{Te}/U)\sqrt{\omega_{He}\omega_{Hi}} \sim 3\cdot10^9\ \mathrm{sec}^{-1}$. The increment of increase is $\sim \sqrt{\omega_{He}\omega_{Hi}} \sim 3\cdot10^8\ \mathrm{sec}^{-1}$, and the wavelength is $\lambda_{11} = 2\pi/K_{max} \sim 2\pi U/\omega_{He} \sim 0.02$ cm ($K\tau_D \sim 1$). The longitudinal wavelength is $\lambda_{11} = 2\pi/K_{11} \sim 2\pi\rho_e \sim 0.2$ cm.

A theoretical estimate of the electron heating time in the case of scattering by these oscillations [11] gives a reasonable order of magnitude: $\tau_e \sim 4\cdot10^{-8}$ sec in the case $H_0 = 800$ Oe, $U \sim 5\cdot10^7$ cm/sec and $T_e \sim 100$ eV. Since this estimate is sensitive to an accurate value of the ratio U/v_{Te} (the change in this ratio changes τ_e by one order of magnitude twice), we must not attribute /31 particular importance to the numerical agreement of the measured and theoretically determined value of τ_e.

The oscillations being considered provide an increase both of the longitudinal and the transverse electron temperatures.

In our experiments, an increase in the temperature is accompanied by a decrease in the amplitude of the whistle magnetic field and the current velocity U. Since a large portion of the whistle energy goes into electron heating in a short period of time (approximately a half period of the oscillations π/Ω), for an accurate determination of the heating time we must solve the Maxwell equation for the whistle field in a plasma, together with Equation (10).

The observed whistle damping time ($\lesssim 2\pi/\Omega$) also agrees with the estimate of the heating time. Actually, the whistle damping time in the case of paired collisions is determined by the relationship [36]

$$\frac{1}{\tau_W} \sim \frac{\nu_{co}}{\omega_{He}} \Omega \left(1 + \frac{K_z^2}{\lambda_{\parallel}^2}\right)^{1/2},$$

where ν_{co} is the frequency of paired electron collisions. Utilizing Formula (14) for the effective frequency of electron collisions, we thus obtain

$$\frac{1}{\tau_W} \sim \frac{u}{v_{Te}} \Omega \left(1 + \frac{K_z^2}{k_{\parallel}^2}\right)^{1/2}. \tag{21}$$

We obtained this same relationship by assuming that the whistle energy goes into electron heating

$$\frac{1}{\tau_W} \equiv \frac{8\pi}{\tilde{H}_z^2} \frac{3}{2} n_0 \frac{dT_e}{dt}.$$

Thus, it is necessary to use (11) for the change in time T_e, and /32 to take the fact into account that for whistles $\tilde{H}_z \sim (U/\Omega)\lambda_{\parallel}H_0$ and $(\lambda_{\parallel}^2 + \lambda_z^2)^{1/2}K_{\parallel} \sim \omega_{pe}^2\Omega/\omega_{He}C^2$. The damping time estimated according to (21) under the experimental conditions ($U \sim 5\cdot10^7$ cm/sec, $v_{Te} \sim 4\cdot10^8$ cm/sec) is $6\cdot10^{-8}$ sec, which practically coincides with the observed whistle damping time (Figure 1b).

Let us now discuss the role of nonmagnetized ion-sound oscillations ($\cos \Theta \sim 1$, $Kz_D \sim 1$). These oscillations may provide isotropization of the electron distribution function and energy transfer with respect to the electron degrees of freedom at the level $W_j/n_0T_e \sim m_e m_i \sim 10^{-3}$ during a period of time which is small as compared with the electron heating time. According to (12) $\tau_i \sim 4 \cdot 10^{-9}$ sec, which is twenty times less than the heating time. For $W_j/n_0T_e \sim 10^{-3}$ the ion-sound oscillations cannot provide the observed electron heating rate: according to (13), $\tau_e^{(j)} \sim 4 \cdot 10^{-7}$ sec, when $U \sim 5 \cdot 10^7$ cm/sec and $T_e \sim 100$ eV, which is five times greater than the observed heating time. (Only when $W_j/n_0T_e \sim 10^{-2}$ can the observed electron heating be due to ion sound.)

For large amplitudes of the high frequency field, the electrons are heated in a period of time which is much less than the time of electron energy loss due to inelastic collisions. Actually, as may be seen from Figure 1d, the electron cooling time is ~ 1 µsec. For small amplitudes of the high frequency field, the electron heating rate decreases greatly ($\tau_e > 10^{-6}$ sec for $\tilde{H}_z \lesssim 200$ Oe), so that the increase in the electron temperature becomes impossible due to strong losses in the case of inelastic collisions. (We should note that the great influence of inelastic collisions on the electron temperature may explain the rapid decrease in T_e with an increase in the neutral gas pressure even for comparatively large values of the amplitude of the high frequency field $\tilde{H}_z \approx 350$ Oe [37].) /33

Let us investigate the results of ion heating. In our experiments, the ion heating is limited when the Larmor radius with the thermal velocity becomes equal to (0.2 — 0.3)α (see Figure 4c). Assuming that the ions obtain energy as a result of the stochastic mechanism of "reduction" of the phase, we obtain the following for the ion heating time from (18):

$$\Delta t = \frac{\omega_{Hi}}{\overline{u}^2}\sqrt{\frac{m_i}{m_e}}\left(\frac{a}{3}\right)^2, \qquad (22)$$

where $\bar{U}$ is the average value of U(t) during the heating time.

It may be assumed that under conditions of excitation resonance (Figure 1), after there is rapid heating of the electrons and the whistle amplitude decreases by a factor of four ($\bar{U} \approx 10^7$ cm/sec); there is stochastic ion heating at the stage of slow damping of a whistle with a small amplitude. The situation is different with departure from excitation resonance toward large magnetic fields H_o, where (Figure 3) the whistle amplitude decreases by approximately a factor of two, and the oscillation damping time in the circuit increases to a value of $\gtrsim$ 1 μsec. The estimate (22) gives $\Delta t \approx (1 — 4)10^{-6}$ sec for these cases, which does not contradict the experimental data (Figure 3).

Paired collisions apparently have a slight influence (even for small $\tilde{H}_z$) on ion heating. Cooling of the ion component may be /34 determined by the charge exchange and (in the case $T_e \gtrsim 5$ eV) by Coulomb collisions with electrons.

We should note that under our conditions when $T_e \gg T_i$ stochastic ion heating in the case of scattering by nonmagnetized ion-sound oscillations ($\cos\Theta \sim 1$), in the case $W_j/n_0T_e \sim 10^{-3}$, is much less intense than for ion-sound oscillations propagated almost perpendicularly to the magnetic field, or for hydrodynamic pulsations ($\tau_z^{(i)} \gg \tau_z$). Nonmagnetized ion-sound oscillations could make a contribution to ion heating only in the case $W_j \sim 10^{-2}n_0T_e$. It must also be noted that this heating mechanism can only operate when $T_e \gg T_i$, whereas stochastic ion heating in the case of ion scattering by hydrodynamic pulsations is in operation when $T_e \lesssim T_i$.

Relationships (17), (18) and (22) were obtained on the assumption that, due to strong interaction of the oscillations, the characteristic correlation time of the Fourier component of the turbulent pulsation potential is on the order of the inverse nonlinear decrement of damping, i.e., on the order of $1/\sqrt{\omega_{He}\omega_{Hi}}$, both for ion-sound oscillations and for hydrodynamic pulsations ($\cos\Theta \sim \sqrt{m_e/m_i}$. This assumption was confirmed by us for hydrodynamic pulsations which were studied in [8, 9]. Using data given in [9], a correlation analysis was made of the measured potential of oscillations $\varphi(t)$ having a maximum of spectral intensity in the frequency region $\omega \approx 1.6 \cdot 10^8$ sec^{-1} $\sim 0.5\sqrt{\omega_{He}\omega_{Hi}}$ in the case $H_0 \approx 900$ Oe and $\bar{n} \approx 10^{13}$ cm^{-3}. The oscillation potential is measured by a double electric probe with a distance between the rods of d = 3 mm. This probe was most sensitive to oscillations with a wavelength ~ d. The perturbation of the plasma caused by the probe (decrease in density, cooling, etc.) was difficult to take into account. The low impedance of the probe in the plasma also distorted the measurements. Thus, the measurements of the /35 oscillation potential $\varphi(t)$ are only qualitative in nature.

Figure 5 shows the normed self-correlation function

$$R(\tau) = \frac{\int_0^T \varphi(t)\varphi(t+\tau)dt}{\int_0^T \varphi^2(t)dt}$$

for one of the realizations of $\varphi(t)$ shown in this same figure, together with the oscillogram of the azimuthal component of the whistle magnetic field $\tilde{H}_\varphi$. The noise oscillograms were studied on a computer. The computer program stipulated subtracting the non-zero average value $\overline{\varphi(t)}$. As follows from Figure 5, the correlation function decreases by a factor of two in the time $\tau \approx 10^{-8}$ sec, which coincides with the opposite value of the oscillation increment of increase. For other realizations of $\varphi(t)$ measured for different values of n, H_0 and $\tilde{H}$, the function $R(\tau)$ had a similar form.

The experiments performed under nonresonance conditions or under resonance conditions, but with a smaller field amplitude ($\tilde{H}_z \sim 500$ Oe) show (see Figures 3 and 4) that with small amplitudes there is primarily ion heating. This is in qualitative agreement with the picture of electron and ion heating examined above (see Section 6). The rapid decrease in the electron temperature with a decrease in the voltage in the circuit, i.e., the whistle amplitude, may be explained by the fact that the rate of electron heating in agreement with (11) rapidly decreases with a decrease in the whistle amplitude. In the case $\text{æ} \gtrsim 1$, the rate at which electrons and ions are heated becomes the same. With a decrease in amplitude, as was already noted above, the role of inelastic electron collisions increases. Since the heating rate is small for small whistle amplitudes, the electron heating generally stops at a certain value of the current velocity $U = U^*$, at which the energy acquired in the field of turbulent pulsations is comparable with the energy lost in the case of inelastic collisions. In this case, the ion heating does not stop, however, since the frequency of inelastic ion collisions with neutral particles is much less than the frequency of inelastic electron collisions. Therefore, predominantly ion heating in the case of small whistle amplitudes in our experiments may be explained by the influence of inelastic collisions which limit the increase in electron temperature. /36

We should also note that for small T_e (under our conditions for T_e = 1 — 10 eV) the analysis of electron heating is complicated by the necessity of considering electron Coulomb collisions with ions, whose presence also leads to bunched-resistance instability [15].

It must be noted that this mechanism of ion heating must not lead to a change in the velocity distribution of ions along the magnetic field. However, the anisotropy of the ion velocity distribution ($T_{\perp i} > T_{\parallel i}$) arising in the case of heating leads to a formation of a different type of anisotropic instabilities.

The reverse influence of the oscillations produced on the ions causes the transfer of transverse ion energy into longitudinal energy. Thus, the isotropization time may be on the same order of magnitude as the heating time. This may possibly explain the fact that in our experiments (see also [7]) the transverse and longitudinal ion energy is identical in order of magnitude throughout the entire heating time.

In conclusion, let us briefly discuss the results of experiments [3, 4] in which there was also turbulent heating of a plasma with a transverse current, caused by a strong high frequency field. /37

In experiments [3] ($\bar{n} \sim 10^{12} - 10^{13}$ cm^{-3}, a = 6 cm, $H_0 \sim H^{\sim} \sim$ 500 Oe, $\Omega = 5.7 \cdot 10^7$ sec^{-1}) a direct magneto-sonic wave with a large amplitude was excited. Under optimal conditions, the absorption of high frequency energy in a plasma takes place in a period of time which is less than the period of oscillations. The electron temperature reached values of 100 — 1000 eV, and the ion temperature — 175 eV. The values of the thermal velocity

$$u = u_\varphi = \frac{c}{4\pi n_0 e}\left|\frac{\partial H_z^{\sim}}{\partial z}\right| \sim \boxed{\frac{\Omega H_z^{\sim}}{4 a n_0 e}} \qquad (23)$$

were rather high ($U_\varphi \approx 3 \cdot 10^8$ cm/sec). Therefore, in these experiments conditions occurred when $U_\varphi \gtrsim v_{Te}$, and Buneman instability took place. Under the same conditions when $U_\varphi < v_{Te}$, ion-sound ($\cos\Theta \sim 1$ or $\cos\Theta \sim U/v_{Te}$) and hydrodynamic ($\cos\Theta \sim \sqrt{m_e/m_i}$) instabilities were excited, which could be responsible for turbulent heating of ions and electrons in the $U_\varphi < v_{Te}$ mode.

In experiments [4] a direct magneto-sonic wave was also excited ($\bar{n} \sim 10^{13}$ cm^{-3}; $T_e \lesssim 10$ eV; $H_0 \approx 2 \cdot 10^3$ Oe; $H^{\sim} \sim 60$ Oe; $\Omega = 1.3 \cdot 10^8$ sec^{-1}). There was thus predominantly ion heating.

The amplitude of the current velocity, according to (23) $U \sim 10^7$ cm/sec was much greater than the thermal ion velocity at the beginning of the heating process and became on the order of $v_{Ti} \sim 10^7$ cm/sec when the ions reached a maximum temperature of $T_i \sim 100$ eV. Therefore, at the initial stage ion heating may occur due to the development of hydrodynamic instability ($\cos\Theta \sim \sqrt{m_e/m_i}$), and due to the small amplitude $\tilde{H}$ electron heating is insignificant. In the case $T_i \sim 100$ eV $\gg T_e$ this instability may change /38 (in the case $U_\varphi \lesssim v_{Ti}$) into electron-sonic instability, which also is accompanied by predominantly ion heating [20 — 23].

The authors would like to thank A. V. Smirnov and I. B. Pinos, who assisted in analyzing the results of measuring ion energy on a computer, and V. T. Pilipenko for their help in the measurements.

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Translated for National Aeronautics and Space Administration under contract No. NASw-2483, by SCITRAN, P. O. Box 5456, Santa Barbara, California, 93108.